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A SOAP OPERA - THE SAD TALE OF THE QUARK

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ABSTRACT

1. INTRODUCTION - QCD, LEADING TERMS AND THE OZI RULE

2. BASIC PHYSICS AND MODELS

  2.1 The Basic Physics of Hadron Structure as Revealed by Experiment

  2.2 The Two Complementary Models

  2.3 Bag Models, Pion Clouds and Multiquark Physics

3. APPLICATIONS OF THE CONSTITUENT QUARK MODEL

  3.1 Hadron Masses

  3.2 Spin Physics

    3.2.1 Baryon magnetic moments

    3.2.2 Some comments on the basic physics of G_A/G_V

    3.2.3 High energy reactions

  3.3 Multiquark Physics

    3.3.1 Bound states - saturation, multiquark clustering and the nucleon-nucleon interaction

    3.3.2 Continuum states - beyond the additive quark model for hadron cross sections

4. THE THEORETICAL BASIS AND PHENOMENOLOGY OF THE OZI RULE

  4.1 Introduction - Non-leading Terms, Nonet Symmetry and Exchange Degeneracy

  4.2 Phenomenology of Charmonium Decays and Three-meson Photoproduction

  4.3 A Topological Criterion for a Forbiddenness Hierarchy

1. INTRODUCTION - QCD, LEADING TERMS AND THE OZI RULE

A SOAP OPERA - THE SAD TALE OF THE QUARK

Alas the poor quark. Condemned to confinement and infra-red slavery for the duration of his normal existence. Freedom comes only in asymptopia - and at the price of destruction of his hadron home by a deep inelastic nuclear explosion.

The soap opera continues after this message:

DID ELLIOTT BLOOM FIND A NEW PARTICLE?

DID ELLIOTT MAKE THE DESERT BLOOM BY FINDING THE BEGINNING OF A NEW PARTICLE PHYSICS BEYOND THE STANDARD MODEL?

To find the answers to these questions: Support your local Physics Department. Lobby for funds for HEP. From your Government. Motivate your best students to go into High Energy Physics instead of genetic engineering. This is the only way to find the answers.

PHYSICS IS GOOD FOR YOU

Keep your eyes on the physics, not on the χ^2 fit. Pay attention to observables of the real world, not to aethers invented by theorists to make themselves feel good. Follow the insight of the MGMJHDJEPW collaboration and the great Nobel Laureate with the initials M. G-M.

MARIA GOEPPERT-MAYER

The shell model marches on. It was obviously wrong but it lead us to new physics.

Now back to the Sad Tale of the Quark. The history of quark physics can be divided into several periods:

1. The Primordial Period. Once upon a time physicists believed that nucleons and pions were elementary particles and that Yukawa's meson theory was analogous to QED for strong interaction dynamics.[1] Then the Δ, ρ and ω were found and it became clear that the nucleon and pion were no longer elementary.

2. The Confused Period of Nuclear Democracy. All particles were then considered equal and connected by a magic bootstrap which explained everything.

3. The Nuclear Age. But there is no magic in the real world and no democracy for particles, only hierarchies. Physicists came to terms with a composite shell model for hadrons made of quarks, analogous to

nuclear physics.[2] Quarks were bound into hadrons with an unknown short-range force whose effective interactions were determined from phenomenological analyses.[3]

4. <u>The Equal Opportunity Model</u>. The recognition of the color degree of freedom as an exact symmetry led to complete freedom and equality for all quarks regardless of color.

5. <u>The "Pie in the Sky when you Die" Model</u>. The introduction of the concepts of Asymptotic Freedom and Infra-red slavery gave freedom to the poor quark only "way up high" in asymptopia. The normal condition is infra-red slavery and confinement within hadrons. The price of freedom is destruction by Deep Inelastic Explosions.

6. <u>The White Supremacy Model</u>. Now there is freedom only for Whites. All colored objects are permanently confined.

7. <u>The Atomic Age</u>. The advent of QCD and the discovery of charmonium led to a picture of hadrons analogous to atoms with a Schroedinger or Dirac equation using a potential obtained from theory.[5]

8. <u>The Condensed Matter Age</u>. Today hadron structure is beginning to look more and more like condensed matter physics. The basic theory is known, the Lagrangian of QCD. But the systems and the dynamical equations are so complicated that nobody knows how to get a solution for a hadron wave function from first principles. As in condensed matter physics, the elementary excitations easily accessible to experiment are described by simple phenomenological models; e.g. potential models, parton models and bag models. The coordinates used for the relevant degrees of freedom need not describe quanta of the fundamental fields, but may be interpreted as quasiparticles or collective coordinates.

How can we use QCD for hadron dynamics when no one knows how to do a full QCD calculation. We have the "leading term" syndrome. Theorists use QCD-motivated hand waving to pick the "leading" contribution, calculate and compare with experiment. If it works they write a paper. If it does not work, they blame "non leading" contributions and write a paper anyway.

How can we attack non-leading effects? There is no clear recipe. We must look for the relevant physics and look for signals in the noise.

The search for charm is an instructive example of theory failure. In their 1974 review, Gaillard, Lee and Rosner[4] instructed experimenters how to look for charm. They suggested everything except looking for very narrow 1^- resonance produced in e^+e^- annihilation. Lipkin[4] suggested a dramatic increase in strange hadron production in e^+e^- annihilation above charm threshold, where second generation c and s quarks would be produced in equal numbers with first generation u and d quarks. Nature fooled Lipkin by cheating and threw in a new player in the game, the τ lepton, which nobody expected and which decayed mainly into nonstrange channels. Thus, both strange and nonstrange production went up and no dramatic increase in strange production was observed.

GLR predicted a 1^- $(c\bar{c})$ state which could be produced in e^+e^- annihilation. But their calculated width was too large by factor 30. Thus, they missed predicting a very striking signal. The reason that GLR went wrong in predicting width of J/ψ is just the non-leading term problem. The hadronic decay of J/ψ is forbidden by the OZI rule. But all this says is that the leading term in the decay amplitude vanishes. However, OZI forbidden decays do occur in nature and must proceed via a non-leading term. This leaves the crucial question of how to estimate the right non-leading term. GLR used the well known OZI-forbidden $\phi\to\rho\pi$ decay to estimate a non-leading term. This turned out to be wrong because the relevant non-leading term $\phi\to\rho\pi$ decay is the two-step transition $\phi\to K\bar{K}\to\rho\pi$, where both transitions are allowed with the $K\bar{K}$ on shell and are related by unitarity to the $\phi\to\rho\pi$ transition. The non-leading term is absent in J/ψ decay where there is no open OZI allowed channel related by unitarity to hadronic decays of the J/ψ.

It is now believed that the J/ψ decay proceeds via a three-gluon intermediate state and that hadronic two-step transitions analogous to $\phi\to K\bar{K}\to\rho\pi$ are suppressed when the intermediate state is off shell. But this was not at all clear at the time of the GLR prediction and there was no previous experimental evidence for OZI forbidden transitions under these conditions. Thus, GLR could not use existing experimental data at that time to estimate the right non-leading term. A more detailed discussion of the OZI rule is given in section 4.

2. BASIC PHYSICS AND MODELS

During the twenty years since quarks were first proposed, an enormous amount of experimental data have been accumulated and interpreted with various versions of the quark model. Each treatment picks a particular "leading approximation." We first review the experimental situation and then discuss the models used to treat them.

2.1 The Basic Physics of Hadron Structure as Revealed by Experiment

2.1.1 Constituent quarks. The low-lying hadron spectrum is described by the states of a quark-antiquark pair for mesons and of three quarks for baryons.[5,6] The relevant degrees of freedom are the spins and flavors of the quarks and/or antiquarks and one relative spatial coordinate for the meson and two for the baryon. There is no evidence for any other degree of freedom in the observed spectrum of low-lying elementary excitations in the 1-2 GeV range. These constituent quarks are clearly not point-like elementary particles but have an effective mass which describes their contribution to the hadron mass and also defines the scale of their magnetic moments.

2.1.2 Identity of mesonic and baryonic constituent quarks. The meson and baryon spectra show that the same constituent quarks appear in both, with the same phenomenological parameters, such as effective masses. Strikingly successful relations between meson and baryon spectra, masses, scattering processes and decays follow from the assumption that both types of hadrons are made from the same quarks.[3,7] These relations do not arise in models which treat mesons and baryons on different footing, like the old Yukawa model or models treating baryons as topological solitons (Skyrmions) and mesons as very different objects.[8]

2.1.3 The universal 200 MeV (1 Fermi) scale. The size of the nucleon and all low-lying nucleon excitations of both parities are all characterized by the same scale which is also the scale of the pion mass. Consider for example, pion photoproduction at low-lying N^* resonances of both parities, where the excitation energies, the mass of the quantum emitted (pion) in the decay and the decay width are all on the same scale as the proton radius. This is in marked contrast to atomic physics and positronium, where the scales of the size of the bound state, the energies of the lowest excitations of the same parity (hyperfine or spin-flip excitations) and the widths of these excited levels and the photon mass all differ by orders of magnitude. The universal scale leads immediately to the paradox that the orbital excitations cannot be described by nonrelativistic motion of a point particle in an orbit with the nucleon radius. Since the energy $\hbar c/r$ corresponding to the radius r and the orbital excitation energies $\hbar\omega$ have the same scale, a point particle moving with frequency ω in an orbit of radius r has a velocity

$$v = r\omega = \left(\frac{\hbar\omega}{\hbar c/r}\right) c \approx c$$

This result has been rigorously derived from the Heisenberg equations of motion.[1)]

2.1.4 Quark number conservation. Although strong interactions can create quark-antiquark pairs, the number of constituent quarks in a hadron appears to be a constant of the motion. There is no experimental evidence for appreciable mixing of states with higher quark numbers in mesons and baryons.

2.1.5 Violation of quark number conservation. Hadrons can emit and absorb mesons. Meson exchange describes an important feature of hadron-handron interactions and in particular of the nucleon-nucleon force relevant to nuclear physics. Thus, both quark number conservation within hadrons and quark-antiquark pair creation in exchanges between hadrons (or meson fields around hadrons) are necessary features of hadron physics.

2.1.6 Saturation. The hadron spectrum saturates at the quark-antiquark and three-quark levels. Even though the quark-antiquark force is known to be attractive in all channels and would be expected to be more strongly bound to three or more quarks than to a single quark, there is very strong experimental evidence against: 1) the existence of strongly bound multiquark states with larger numbers of constituents like a dipion state with a mass less than the mass of two pions; 2) πN reactions analogous to the common nuclear stripping reaction $t(d,p)\alpha$, in which an antiquark is stripped from the pion and bound to the three quarks in the nucleon, $p(\pi,q)\ qqq\bar{q}$.

2.1.7 Multiquark clustering in a lower 8 MeV energy scale. Many multiquark states are observed as nuclei, bound by energies much lower than the hadronic scale and described to a good approximation as assemblies of nucleons. Their structure at the quark level is an assembly of three-quark clusters. This suggests a picture with hadrons behaving in QCD-like neutral atoms, and forces between hadrons much weaker than those binding quarks into hadrons. Nuclei are thus analogous to molecules and the residual forces binding quark-clusters into nuclei are analogous to molecular forces.

2.1.8 Broken-string confinement. Isolated quarks have never been observed. However, there is no experimental evidence for a so-called "confining potential" between constituent quarks which rises to infinity and prevents the separation of the constituents. The basic physics is described more properly by a string model for hadrons with quarks at the ends. Pulling constituent quarks apart does not require infinite energy; it creates mesons by breaking the string.

2.1.9 Quark additivity. A large variety of processes are described by the additive quark model (AQM) in which a single quark in a hadron is active in all transitions and the remainder are spectators.[9] These include electromagnetic, weak semileptonic and strong mesonic decays of hadrons and hadron-hadron total cross sections and scattering and reactions at low momentum transfers.

2.1.10 Simple hadron mass formulas. Hadron masses are well described by simple shell models in which the hadron mass is the sum of the constituent quark masses and an effective two-body interaction with the color, spin and flavor dependence of a one-gluon exchange potential in QCD.

2.1.11 Current quarks. The weak and electromagnetic currents appear to couple to point-like objects with quark and antiquark quantum numbers in deep inelastic scattering; i.e. in processes with high momentum transfer. The algebra of these currents seems also to describe some low-energy properties of hadrons. However, these current quarks are not the same as the constituent quarks, and the hadron wave functions with current quarks are not simply given by any model.

2.1.12 The multiquark continuum. Hadron-hadron scattering is described by phenomenological short-range interactions. In principle, these should be obtainable from an underlying quark description.

2.2 The Two Complementary Models

With these basic physical points in mind, we can see the phenomenology of the two basic and complementary models, the Quasinuclear Colored Quark Model (QNCQM) and the quark-parton model (QPM). The QNCQM uses constituent quarks and model wave functions given by some phenomenological potential model. But the constituents have completely unknown properties (except for the conserved quantum numbers). The QPM has well-known and completely defined point-like constituents in a completely unknown hadron wave function. In the QNCQM the properties of the quarks are free phenomenological parameters; e.g. masses and form factors, which are determined from experiment. In the QPM the properties of the hadron wave function are free phenomenological parameters (structure functions) which are determined from experiment. In both cases a large amount of experimental data are available on the relevant elementary excitations, and significant fits to the data and new predictions subsequently verified by new data have been made after the parameters were determined.

In the modern version of the QNCQM[1,10,11] with a two-body color-exchange force a number of the basic physical features of hadron structure listed above arise automatically, whereas others do not. We list these in detail.

2.2.1 The constituent quark description of the low-lying spectrum comes out of the model directly.[12]

2.2.2 A universal description of mesons and baryons with the same constituents and the same Hamiltonian arises naturally in the model.

2.2.3. Saturation and the absence of deeply bound multiquark states come naturally from the color-exchange force.[10,11]

2.2.4 Simple mass relations arise from the additional assumption of the Fermi hyperfine interaction due to one-gluon exchange,[3,5,13] again relating mesons and baryons and giving new predictions for baryon magnetic moments.

2.2.5 The addition of the quark additivity assumption gives further successful predictions for a large variety of processes.

2.2.6. Quark number conservation and violation respectively in hadron wave functions and decays do not arise naturally in the model. Put in by hand they lead to successful predictions respectively for the spectrum and for transitions again showing a universal description of mesons and baryons.

2.2.7 Multiquark clustering has recently been described by this same model, giving a deuteron-like molecular bound state[14-15] and a reasonable description of the deuteron.

2.2.8 Broken-string confinement is not yet described in any simple way.

2.2.9 The single scale and its implications for relativistic motion do not fit together with this model in any consistent way.

2.2.10 Current quarks. The model has no pretensions for describing deep inelastic scattering, although there have been attempts to introduce a structure for the constituent quarks in terms of current quarks for describing processes at high momentum transfer.

2.2.11 The multiquark continuum. Hadron-hadron scattering is not easily described by models with constituent quarks and two-body interactions, because of the unphysical long-range Van-der-Waals forces that arise in these models.

The quark-parton model can be considered as a more fundamental description with the well-defined constituents as given by QCD. However, it suffers from providing no method for a feasible calculation of hadron wave functions and low energy properties. Perhaps lattice gauge theories will prove to be the answer. But so far they have not given the solution.

The constituent quark picture can be justified by hand-waving arguments in which the valence quarks of the parton model each acquire a share of the gluons and ocean of quark-antiquark pairs present in the wave function and define a quasi-particle with a finite mass and size. Open questions are: 1) why these additional degrees of freedom can be absorbed into the constituent quarks for low energy spectroscopy and leave no residual additional degrees of freedom; 2) why constituent quarks can be described by nonrelativistic quantum mechanics, in view of eq. (2.1). Some relativistic corrections have been shown to be appreciable but absorbed into the definitions of the effective masses of the quasiparticles. Such corrections simply renormalize the effective mass parameters and are automatically included if the parameters are adjusted to fit experiment.[16] But no convincing overall justification exists except for the excellent agreement with experiment indicating a choice of the right degrees of freedom to describe the elementary excitations.

2.3 Bag Models, Pion Clouds and Multiquark Physics

Between these two complementary and very successful models for orthogonal physical properties is the bag model.[17] Intuitively this treats the valence quarks of the parton model as point-like zero-mass quarks in a fully relativistic description, and "sweeps the remaining degrees of freedom into a bag." This model has the desirable feature of introducing confinement in a relativistic framework and enabling fully relativistic calculations of hadron properties It also suggests a phase transition in the vacuum in the presence of the quarks and the color field which confines the color field. However, there is no experimental evidence so far for the phase transition, for the presence of bag degrees of freedom in the elementary excitations, nor for new physics revealed by the use of the fully relativistic calculations and not already present in the nonrelativistic constituent quark model. Bag model calculations have generally attempted to treat the same hadron properties already treated successfully by the nonrelativistic model, and have not obtained better results with any significance. They have not obtained new results for properties not given by the constituent model, and predictions of new properties of multiquark systems resulting from the bag have generally failed to agree with experiment.[18]

One interesting feature of the bag model which may be relevant for nuclear physics is the introduction of the pion cloud in the so-called "cloudy bag models."[19-20] Here it may be possible to attack a different part of the physics which is not easily handled by the other models. However, there seems to be a bit of a wild goose chase in the calculations of baryon magnetic moments and of G_A/G_V where it is difficult to obtain a significant improvement over the constituent quark model and any signal is lost in the noise.

There has been considerable confusion about the application of the bag model to multiquark systems. An essential weakness of the model for treating bound multiquark states like complex nuclei is its inability to describe correlations and clustering in multihadron states.[18] The deuteron, for example, is clearly a six quark system which behaves like two three-quark clusters a great deal of the time. In a bag picture it must be described at different times as two separate bags, two overlapping bags, and a single deformed bag with a variety of shapes. There is no simple treatment of the transitions between these different descriptions. The constituent model can treat the six quark system with phenomenological two-body forces and derives the clustering from the dynamics.[15]

Because confinement is put in by hand in the bag model, rather than deriving it from a Hamiltonian as in nonrelativistic potential models, some physical intuition is needed in the application to multiquark systems. These systems are not confined; they can break up into two or more hadrons. Thus, states obtained from solving a multiquark problem in a bag should not be interpreted as physical hadron resonances. The coupling of these states to the open breakup channels is crucial and must be included. Jaffe and Low[21] have developed a formalism for treating this problem, similar to the formalism used in nuclear physics to treat nuclear resonances which can decay by breakup.

3. APPLICATIONS OF THE CONSTITUENT QUARK MODEL

We now consider several applications of the constituent quark model.

3.1 Hadron Masses

Hadron masses have been fit very well by an extremely simple shell model Hamiltonian with a single-particle term called an effective mass and a residual two-body interaction,

$$H = \sum_i m_i + \sum_{i>j} \frac{\sigma_i \cdot \sigma_j}{m_i m_j} u_{ij} \tag{3.1}$$

where m_i and σ_i denote the effective mass and spin of constituent quark i, and u_{ij} is an effective interaction. The original versions of this

Hamiltonian had more general two-body interactions and obtained results relating baryon masses.[3,4] It was developed independently in 1966 at the Weizmann Institute and by Sakharov and Zeldovich in Moscow. The assumption that all the flavor dependence in the residual two-body interaction was in a spin-dependent hyperfine interaction led to a remarkable relation between meson and baryon masses,[3,7]

$$M_{\Lambda} - M_N = m_s - m_d = (3/4)(M_{K^*} - M_{\rho}) + (1/4)(M_K = M_{\pi}) \tag{3.2a}$$

The left hand side of eq. (3.2a) is 177 MeV; the right hand side is 180 MeV, giving striking support to the assumption that mesons and baryons are made of the same quarks.

I discovered this relation in 1978, using the ideas of QCD for the spin dependence of the force.[5] I later learned that it had already been obtained in 1966 by Sakharov and Zeldovich, when I received a post card from Sakharov.[22] He was very kind in quoting the formula (3.2a) with the comment "Of course you are right," rather than angrily pointing out that he had published the same formula twelve years earlier.

The fascinating story of my correspondence with Sakharov is the subject of another talk.[22] The first post card was featured in an editorial in the Washington Post entitled "A Voice out of the Darkness" describing the breakup of the original Moscow seminar on collective phenomena by the KGB. It was reproduced in many newspapers and journals and the story was broadcast in Russian over Voice of America where it eventually reached Sakharov. He answered with another post card, shown in ref. 22).

The Hamiltonian (3.1) gives[3] another relation between meson, baryon and quark masses,

$$\frac{M_{\Delta}-M_N}{M_{\Sigma^*}-M_{\Sigma}} = \frac{M_{\rho}-M_{\pi}}{M_{K^*}-M_K} = \frac{m_s}{m_u} \tag{3.2b}$$

The ratio of baryon mass differences is 1.53; the ratio of meson mass differnces is 1.61. The small discrepancy was later explained by a refinement of a model involving the different sizes of baryon and meson wave functions.[16] The relation with effective quark masses follows from the assumption that the spin dependent interaction is a one-gluon-exchange hyperfine interaction inversely proportional to the constituent quark masses.[5] An equivalent formula for the quark mass ratios appears in Sakharov's second post card, along with the comment that the masses are of course effective masses.

3.2 Spin Physics

3.2.1 Baryon magnetic moments. Baryon magnetic moments have been fit remarkably well by extremely simple constituent quark models. The original broken-SU(6) model[23] assumed SU(6) wave functions for the baryons and used the proton and Λ moments as input to determine the nonstrange and strange quark moments and predict the others. The assumption that the quarks have Dirac moments with the same effective mass parameters m_i as in the Hamiltonian (3.1) led to two successful predictions of the Λ moment[5,24] at the 1% level from eqs. (3.3) and 3.2) respectively.

A determination of all the quark mass parameters from the Hamiltonian (3.1) and the baryon masses enables a prediction of the baryon magnetic moments with no free parameters.[25] Table 1 shows a comparison of recent experimental data[26-29] with Rosner's 1980 predictions[25] and those of the original broken-SU(6) model with two free parameters. The basic question is whether the small discrepancies between experiment and these simple models at the 15% or 20% level arise from a signal or from noise; i.e. from a single dominant dynamical mechanism or from many independent mechanisms of roughly equal strength. It is very difficult to get a better fit with models which must fine tune parameters and thereby lose the beautiful connection between the physics of magnetic moments and hadron masses.

Cloudy bag models therefore encounter difficulty in obtaining significant fits. This is seen in the predictions of one cloudy bag calculation[20] in Table 1. A similar inconclusive result has been obtained by adjusting the strengths of various contributions to fit the data.[19] The better fit obtained is not worth the price of introducing additional free parameters, losing all relations with hadron masses and using physically unreasonable values for quark contributions; i.e. setting the effective magnetic moment of the strange quark to be larger in magnitude than that of the down quark. The pion field may well be a missing ingredient in the simple magnetic moment calculations. But it is not obviously the dominant missing ingredient. There is little point in trying different χ^2 fits to noise if there is no clear signal present. Only when enough other hadron physics not treated by the simple model convincingly demonstrates a consistent method of adding the pion physics to constituent quarks can there be a significant attack on the baryon moments with all parameters fixed by other data, as is the case with constituent quarks.

Table 1. Experimental Values and Model Predictions for Baryon Magnetic Moments

Baryon Moment	1981 Value Refs[25,27,28]	1983 Data Ref[26]	From Naive Model[25]	From Broken SU(3)[23]	From Cloudy Bag[20]
$\mu(p)$	2.793±0.000	2.793±0.000	2.79	*	2.60
$\mu(n)$	-1.913±0.000	-1.913±0.000	-1.86	-1.86	-2.01
$\mu(\Lambda)$	-0.613±0.005	-0.613±0.005	-0.58	*	-0.58
$\mu(\Sigma^+)$	2.33±0.13	2.38±0.02	2.68	2.68	2.34
$\mu(\Sigma^-)$	-0.89±0.14[29]	-1.11±0.04[27]	-1.05	-1.04	-1.08
$\mu(\Xi^\circ)$	-1.25±0.014	-1.25±0.014	-1.40	-1.43	-1.27
$\mu(\Xi^-)$	-0.75±0.06	-0.60±0.04	-0.47	-0.50	-0.51

An alternative approach attempts to analyze the small discrepancies between experiment and the predictions of the simple model without theoretical prejudices to see whether the systematics can give clues to the existence of a relevant degree of freedom.[28] Of particular interest are the deviations of the Σ^+ and Ξ^- moments from the predicted values, and in particular the differences between these moments and the proton and Λ moments respectively. Both these differences are predicted to vanish in the SU(3)-breaking mechanism of increasing the strange quark mass and thus decreasing the magnitude of its moment relative to the nonstrange quarks.[23] The Σ^+ - proton difference is too large to be explained by this mechanism, and the Λ - Ξ^- difference has the wrong sign.

The Σ^+ and proton moments are dominated by their u-quark contributions which are predicted to be equal. The observed Σ^+-p difference is much too large to be explained by the difference between the small s and d quark contributions. This indicates that the nonstrange contributions in the two baryons must be different, either enhanced in the nucleon or quenched in the hyperon. A pion cloud contribution gives precisely this nonstrange enhancement in the nucleon and naturally gives an improvement in the value of the Σ^+ moment[20] as shown in Table 1. However, the additional isovector contribution from the pion cloud destroys the balance between isoscalar and isovector contributions in the nucleon which gave one of the first strikingly successful quark model results; namely -3/2 for the ratio of the proton and neutron moments. This is also seen in Table 1.

3.2.2 Some comments on the basic physics of G_A/G_V. Neither of the two complementary models can predict G_A/G_V. The constituent quark model can only predict this ratio for the nucleon if the value for the quark is known. In the quark-parton model with current quarks $G_A/G_V = 1$, but there is no reason to assume this for constituent quarks. the erroneous value of 5/3 for G_A/G_V of the nucleon is only obtained by making the erroneous assumption that $G_A/G_V = 1$ for constituent quarks. The experimental value of 1.2 is easily fit by taking G_A/G_V for the quark as a free parameter and adjusting its value to be (3/5)1.2.

Bag model calculations take three current quarks and obtain a value below 5/3 by including the relativistic corrections from the small components of the wave function. But most other corrections to the three-valence-quark wave function also reduce G_A/G_V below the value 5/3. Relativistic corrections are not obviously more important than other effects like configuration mixing, pion clouds, etc. With only one number to fit and a variety of corrections each with free parameters available, it is very difficult to get a significant fit and show that the particular model used has the right physics.

The real physics in G_A/G_V was given by Adler and Weisberger using current quarks and current algebra in a consistent description with unknown nucleon wave functions. The necessary information on the proton wave function was taken from experiment; namely in the matrix elements of the axial current between the nucleon states and the pion-nucleon continuum, as given by PCAC and scattering data. Saturating their sum rule by including only the nucleon state gives $G_A/G_V = 1$, corresponding to an elementary fermion; i.e., a quark, which has no elementary excitations induced by the axial current. Saturating the sum rule with the nucleon and the Δ gives $G_A/G_V = 5/3$ corresponding to a system of three current quarks with only spin-flip elementary excitations induced by the axial current.

The Adler-Weisberger result shows an intimate connection between G_A/G_V and the matrix elements of the axial current between the nucleon and the pion-nucleon continuum which can be considered as dominated by higher nucleon resonances. This connection presents an open challenge to any model with current-quark wave functions; namely to calculate not only the value of G_A/G_V but also the contributions of the higher baryon resonances to the Adler-Weisberger sum rule and show that these agree with experiment.

3.2.3 High energy reactions. Both the baryon magnetic moment predictions and the predictions of G_A/G_V depend upon the spin-isospin structure of the baryon wave functions. Another possible experimental test of this structure has recently been suggested,[29,30] following the observation that the process $\pi^- p \to \rho^- p$ seems to be dominated by a helicity flip.[31]

The basic physics is in the structure of the proton wave function where the fermi statistics of colored-quarks in a color singlet state require the two u-quarks to be antisymmetric in color and therefore symmetric in space and spin. Since the dominant part of the wave function is spatially symmetric, the spins of the two u-quarks are parallel and the d-quark spin is antiparallel to the u-quark spin to give a total spin of 1/2. Thus the dominant piece of a proton wave function with positive helicity has both u-quarks with positive helicity and the d-quark with negative helicity.

In any model for a helicity flip transition where the d-quark is active and the u-quarks are spectators, this dominant part of the wave fuction cannot contribute to the transition amplitude. Flipping the helicity of the d-quark produces a three quark state with all spins parallel and which has no overlap with the proton. The transition must go via a smaller component in the wave function where the d-quark has positive helicity. Examples of transitions where the d-quarks in the proton is active and the u-quarks are spectators are: 1) The reaction $\pi^- p \to \rho^- p$ via constituent exchange — the d-quark in the π^- exchanges with the d-quark in the proton and the u-quarks are specators. 2) The reaction $\pi^+ p \to \rho^+ p$ via $q\bar{q}$ annihilation. The $\bar{d}$ antiquark in the π^+ must annihilate against the d-quark in the proton with subsequent $d\bar{d}$ pair creation, and the u-quarks are spectators.

In both these models the helicity flip transitions for the "mirror" reactions $\pi^- p \to \rho^- p$ and $\pi^+ p \to \rho^+ p$ are expected to be very different, with the $\pi^- p \to \rho^- p$ strongly suppressed in quark exchange and $\pi^+ p \to \rho^+ p$ strongly suppressed in annihilation. Comparison of the two reactions can therefore provide a sensitive probe of the reaction mechanisms.

This argument can be stated quantitatively by evaluating transition matrix elements with the use of the Wigner-Eckart theorem in both spin and isospin spaces.

$$\frac{\langle p\uparrow|S_{u+}-S_{d+}|p\downarrow\rangle}{\langle p\uparrow|S_{u+}+S_{d+}|p\downarrow\rangle} = \frac{\langle p\uparrow|(\vec{S}_u-\vec{S}_d)\cdot(\vec{S}_u+\vec{S}_d)|p\uparrow\rangle}{\langle p\uparrow|(\vec{S}_u+\vec{S}_d)\cdot(\vec{S}_u+\vec{S}_d)|p\uparrow\rangle} = \frac{S_u(S_u+1)-S_d(S_d+1)}{S(S+1)} = \frac{5}{3} \quad (3.3a)$$

$$\frac{\langle p\uparrow|S_{u+}|p\downarrow\rangle}{\langle p\uparrow|S_{d+}|p\downarrow\rangle} = \frac{\langle p\uparrow|S_{uz}|p\uparrow\rangle}{\langle p\uparrow|S_{dz}|p\uparrow\rangle} = \frac{\langle p\uparrow|S_{uz}|p\uparrow\rangle}{\langle n\uparrow|S_{uz}|n\uparrow\rangle} =$$
$$= \frac{\langle n\uparrow|S_{dz}|n\uparrow\rangle}{\langle n\uparrow|S_{uz}|n\uparrow\rangle} = \frac{S(S+1)+S_u(S_u+1)-S_d(S_d+1)}{S(S+1)-S_u(S_u+1)+S_d(S_d+1)} = -4, \quad (3.3b)$$

where $\vec{S}_u$, and $\vec{S}_d$ are the total spin operators for the u quarks and d quarks respectively in the nucleon and $\vec{S}=\vec{S}_u+\vec{S}_d$.

This result depends on an assumed proton wave function only in the set of values $S_u=1$, $S_d=1/2$ and $S=1/2$ introduced to give the final equalities and the values 5/3 and −4. These values can be seen to follow from very general considerations. Any wave function with only one valance d-quark has $S_d=1/2$. Any wave function with zero orbital angular momentum has $S=J=1/2$. Any wave function with two valence u-quarks, an antisymmetric color coupling to give a color singlet and even value for the relative orbital angular momentum betwen the two identical u-quarks is required by Fermi statistics to have $S_u=1$. An explicit derivation of the conditions on the proton wave function needed to obtain the result (3) is given below.

The relation (3) has given the successful prediction of the ratio of the proton and neutron magnetic moments

$$\frac{\mu_p}{\mu_n} = \frac{\langle p\uparrow|2S_{uz}-S_{dz}|p\uparrow\rangle}{\langle n\uparrow|2S_{uz}-S_{dz}|n\uparrow\rangle} = \frac{2\langle p\uparrow|S_{uz}|p\uparrow\rangle/\langle p\uparrow|S_{dz}|p\uparrow\rangle - 1}{2 - \langle n\uparrow|S_{dz}|n\uparrow\rangle/\langle n\uparrow S_{uz}|n\uparrow\rangle} = -\frac{3}{2} \qquad (3.3c)$$

where the factors -2 comes from the ratio of the electric charges of the u and d quarks.

We now apply the results (3.3b) to the helicity flip processes $\pi^- + p\downarrow \to \rho^- + p\uparrow$ and $\pi^+ + p\downarrow \to \rho^+ + p\uparrow$

If these processes proceed via some combination of gluon exchanges, the gluons do not distinguish between quark flavors and the cross sections for the two processes are equal

$$\left.\frac{\sigma(\pi^- p\downarrow\to\rho^- p\uparrow)}{\sigma(\pi^+ p\downarrow\to\rho^+ p\uparrow)}\right]_{\text{Gluon}} = 1 \qquad (3.4a)$$

If, however these processes proceed via quark exchange, then the reaction $\pi^-\to\rho^-$ involves the interchange of a d-quark in the π^- with a d-quark in the proton; the $\pi^+\to\rho^+$ reaction involves the interchange of a u-quark in the π^+ with a u-quark in the proton. In all cases the exchanged quark is replaced by another quark of the same flavor and opposite helicity and the transition matrix elements are proportional to the matrix elements of the spin-flip operators S_{u+}, S_{u-}, S_{d+} and S_{d-}. By isospin symmetry

$$\langle\rho_\alpha^-|S_{ui}|\pi^-\rangle = \langle\rho_\alpha^+|S_{di}|\pi^+\rangle$$

$$\langle\rho_\alpha^-|S_{di}|\pi^-\rangle = \langle\rho_\alpha^+|S_{ui}|\pi^+\rangle$$

for any component i of the operators S_u and S_d and any polarization state α of the ρ. Thus the transitions differ only in the baryon transitions, and

$$\frac{\sigma(\pi^- p\downarrow \to \rho^- p\uparrow)}{\sigma(\pi^- p\downarrow \to \rho^+ p\uparrow)}\Big]_{QEX} = \left| \frac{\langle p\uparrow | S_{d+} | p\downarrow\rangle}{\langle p\uparrow | S_{u+} | p\downarrow\rangle} \right|^2 = \frac{1}{16} \tag{3.4b}$$

On the other hand, if the process proceeds via quark-antiquark annihilation, the $\pi^- \to \rho^-$ reaction involves the annihilation of the $\bar{u}$ antiquark in the π^- against a u-quark in the proton and the subsequent creation of a $u\bar{u}$ pair, the $\pi^+ \to \rho^+$ reaction involves the annihilation of the $\bar{d}$ antiquark in the π^+ against the d-quark in the proton and the subsequent creation of a $d\bar{d}$ pair.

Again the transition matrix elements are proportional to spin-flip matrix elements. The meson transitions are the same for the two reactions while the ratios of the baryon transition involves rations of spin-flip matrix elements. However in this case the ratio is reversed:

$$\frac{\sigma(\pi^- p\downarrow \to \rho^- p\uparrow)}{\sigma(\pi^+ p\downarrow \to \rho^+ p\uparrow)}\Big]_{ANN} = \left| \frac{\langle p\uparrow | S_{u+} | p\downarrow\rangle}{\langle p\uparrow | S_{d+} | p\downarrow\rangle} \right|^2 = 16 \tag{3.4c}$$

It can be shown explicitly that these results (3.4) hold for any wave function with three valence quarks in a totally symmetric spatial state and a sea of gluons and quark-antiquark pairs with vacuum quantum numbers.

The striking difference between the values (3.4a), (3.4b) and (3.4c) suggest that measurements of the two very similar reactions can provide interesting insight into the underlying dynamics.

3.3 Multiquark Physics

3.3.1 Bound states - saturation, multiquark clustering and the nucleon-nucleon interaction. The QNCQM introduces the color degree of freedom into the shell-model Hamiltonian, (3.1) and multiplies the residual interaction by a color factor F_{ij} which is assumed to depend on the color state of the i-j pair like one-gluon exchange. A spin-independent interaction v_{ij} and the kinetic energies t_i of the quarks are also included.

$$H_{QNCQM} = \sum_i (m_i + t_i) + \sum_{i>j} F_{ij}\left(v_{ij} + \frac{\sigma_i \cdot \sigma_j}{m_i m_j} u_{ij}\right) \tag{3.5}$$

The color dependence of F_{ij} is not justified from first principles for an effective interaction which should also include higher order exchanges. However, the experimental agreement with this assumed color dependence is very striking.[1,11] The Hamiltonian (3.5) is assumed to hold for any system of quarks and antiquarks with the same parameters. It immediately gives the result that only quark-antiquark and three-quark systems are strongly bound[11] and leads to relations between meson and baryon spectra[7] in agreement with experiment. This Hamiltonian has also been used to derive the simpler shell model Hamiltonian (3.1) by showing that both the kinetic energy terms t_i and the effective matrix elements of the spin-independent interaction v_{ij} can be absorbed [16] in the effective mass terms m_i to a good approximation. This point has also been noted by Sakharov.[22,32]

One very interesting problem in multiquark physics is to obtain a more fundamental understanding of the nucleon-nucleon interaction from the quark picture. The principal physical features and difficulties of this problem have been recently clarified. The short-range repulsion and the intermediate range attractions of the nucleon-nucleon force can be obtained from a quark model for the two-nucleon system and a reasonable quark-quark interaction.[15]

The problem, however, is complex, and has many pitfalls. Errors arise from improper treatment of:

1. Broken-string confinement. Confining potentials which increase to infinity at large distances give rise to spurious long range Van-der-Waals forces.[1,33]
2. Pauli blocking. An elementary neutron and an elementary proton with parallel spins can be at the same point in space without violating the Pauli principle. Three u quarks and three d quarks can also be at the same point without violating the Pauli principle. The dominant components in the six-quark wave function for a neutron and proton with spin up can also be at the same point in space without violating the Pauli principle; namely two u-quarks with spin up and one d-quark with spin down for the proton $|u\uparrow u\uparrow d\downarrow\rangle$, and two d-quarks with spin up and one u-quark with spin down for the neutron $|d\uparrow d\uparrow u\downarrow\rangle$. But these components constitute only 2/3 of the wave function of each nucleon and 4/9 of the

two-nucleon wave function.

The remaining 1/3 of the proton wave function, $|u{\uparrow}u{\downarrow}d{\uparrow}\rangle$, and the remaining 1/3 of the neutron wave function, $|d{\uparrow}d{\downarrow}u{\uparrow}\rangle$, each encounter Pauli blocking when combined with the dominant component of the other nucleon. They lead to the six-quark states

$$|(u{\uparrow}u{\downarrow}d{\uparrow})_1;(d{\uparrow}d{\uparrow}u{\downarrow})_1\rangle \tag{3.6a}$$

$$|(u{\uparrow}u{\uparrow}d{\downarrow})_1;(d{\uparrow}d{\downarrow}u{\uparrow})_1\rangle \tag{3.6b}$$

where the subscript 1 denotes that the three-quark states in parenthesis are coupled to a color singlet. The state (3.6a) contains three d quarks with spin up. These can be at the same point in space only if they are totally antisymmetric in color and are therefore in a Δ^- state. Once the quarks are pushed together to the same point this six-quark state is required to have the color and spin couplings of a Δ - N or Δ - Δ system.

Another way to see this problem is to note that the Pauli principle for the $d{\uparrow}$ quarks in the wave function (3.6a) requires the three to have different colors. But a two-nucleon wave function has no color-correlations between the quarks in the proton and the quarks in the neutron. The probability that the three $d{\uparrow}$ quarks in the wave function (3.6a) have the right colors to be consistent with the Pauli principle is 1/3. A similar argument for the two $u{\downarrow}$ quarks in the wave function (3.6a) gives and additional factor of 1/2. Thus the wave functions (3.5) with each nucleon coupled to a color singlet has a probability of 5/6 of violating the Pauli principle when the two nucleons are brought together.

A similar argument holds for the state (3.6b). Since each of these states constitutes 2/9 of the total two-nucleon wave function and each has a probability of 5/6 of violating the Pauli principle when the two nucleons are brought together, we see that 10/27 or 37% of the two-nucleon wave function violates the Pauli principle.

3. The Heisenberg uncertainty principle. Additional kinetic energy arises when two nucleons are brought together because of the changes in

the wave function at small distances. The Pauli blocking seen in the states (3.6) introduces spatially dependent color-spin couplings and spatial derivatives into the wave function. These increase kinetic energy. This effect can be also seen by noting that the Pauli blocking introduces a Δ in the wave function at small distances. But confining the Δ to a small volume costs kinetic energy which does not appear as kinetic energy of the two-nucleon system.

4. The interplay of Fermi, Pauli and Heisenberg. The appearance of the Δ in the six quark wave function required by Pauli at short distance costs Fermi hyperfine energy, since the color and spin couplings of this state give a higher hyperfine energy than that of a two-nucleon state. This can be checked by explicit calculations of the hyperfine energy for the wave functions (3.6) using the color-spin algebra introducted by Jaffe.[13] The potential energy is minimized by confining the Δ color-spin correlations to the minimum volume required by the Pauli principle. But such spacial confinement costs kinetic energy due to the uncertainty principle.

Maltman and Isgur[15] have shown the above features in a calculation using variational wave functions and the Hamiltoniam (3.5) for the six-quark system with simplified forms for the interactions v_{ij} and u_{ij} previously used successfully for baryon spectroscopy.[6] They show that clustering into two three-quark nucleons naturally arises in choosing the wave function to minimize the energy, that an intermediate range attractive force arises from the analog of molecular Van-de-Waals forces and that the interplay of Pauli, Fermi and Heisenberg described above gives a repulsive core. A large space of trial wave functions is needed in order to include all the essential physics.

However, the model has unphysical harmonic confining forces and long-range Van der Waals forces. The authors contend that their calculations and results are insensitive to these long-range effects. But they have yet to show convincingly that their intermediate-range attraction will survive if the infinitely rising quark-quark potential is cut off at a reasonable radius.

The same approach[14] suggests that the δ and S^* scalar meson resonances are deuteron-like $K - \bar{K}$ bound states. Tests of this picture have been proposed using the production of these mesons on complex nuclei.[18] This would be an interesting experiment at the intersection of nuclear and particle physics.

3.3.2 <u>Continuum states - beyond the additive quark model for hadron cross sections</u>. The additive quark model has been remarkably successful in predicting "leading term" relations betwen hadron scattering processes. In one attempt to go beyond the leading term analyses of the small disrepancies between experiment and predictions for hadron total cross sections have revealed striking systematics in the data and led to empirical relations with great and successful predictive power.[28] A clear signal above the noise level was found in the relation and energy dependence of two quantities which are predicted to vanish in the naive SU(3) symmetric additive quark model,[34]

$$\sigma(\pi^- p) - \sigma(K^- p) = (1/3)\sigma(pp) - (1/2)\sigma(K^+ p) \approx B(P/20)^{-\varepsilon} \qquad (3.7)$$

where P denotes the laboratory momentum in GeV and B and ε are parameters.

The two differences appearing in eq. (3.7) represent the deviations from SU(3) symmetry and quark model additivity respectively, and were not expected to be related. They are related in a non-additive "Two-Component Pomeron model"[34] which breaks both additivity and SU(3) symmetry by introducing a double exchange contribution, parameterized by adding an ad hoc non-additive two-body term to the Pomeron contribution, without requiring any SU(3) breaking in the additive component. The equality (3.7) holds over a wide energy range and is fit by a monotonically decreasing function (3.7) with the parameters $B \approx 4.4$ and $\varepsilon \approx 0.2$.

It was then found that the complicated energy dependence of $\sigma_{tot}(pp)$ could be simply expressed as the sum of an SU(3) symmetric monotonically rising component and the decreasing component (3.7). This led to a universal formula for the total cross section for any hadron H on a proton,

$$\sigma_{tot}(HP) = AN_q(H)(P/20)^{\delta} + BN_q(H)N_{ns}(H)(P/20)^{-\epsilon} +$$

$$+C\{N_{\bar{d}}(H) + 2N_{\bar{u}}(H)\}(P/20)^{-1/2} \qquad (3.8)$$

where $N_q(H)$, $N_{ns}(H)$, $N_{\bar{d}}(H)$ and $N_{\bar{u}}(H)$ denote respectively the total number of quarks and antiquarks, the number of nonstrange quarks and antiquarks, the number of $\bar{d}$ antiquarks in hadron H. the first two terms on the right hand side of (3.8) describe the two components of the Pomeron contribution. The third term does not contribute to proton-proton scattering and expresses a Reggeon exchange contribution described by duality diagrams for hadrons containing nonstrange antiquarks.

This model has successfully predicted new experimental results at higher energies with no adjustment of parameters.[35] Its most recent success is in the predictions for hyperon-nucleon total cross sections.[30-38]

$$\sigma(pp) - \sigma(\Sigma p) = \sigma(\Sigma p) - \sigma(\Xi p) = (3/2\{\sigma(\pi^- p) - \sigma(K^- p)\} \qquad (3.9)$$

The AQM[9] assumes SU(3) symmetry breaking at the quark level by an ad hoc difference between strange and nonstrange quark contributions chosen to fit experiment and predicts all differences in eq. (3.9) to be equal; i.e, without the factor 3/2. The new data confirm the factor 3/2 in the prediction (3.9) and are in strong disagreement with the AQM.

This two-component Pomeron model pinpoints certain features of the experimental data which have a simple physical interpretation. The failure of the AQM in relations like (3.9) between the meson and baryon sectors can be attributed entirely to the contributions from nonstrange quarks. This is most clearly demonstrated by projecting out the contributions of strange and nonstrange quarks from the experimental baryon-nucleon and meson-nucleon cross sections.

$$\sigma(nN)_B = (1/6)\{\sigma(pp) + \sigma(pn)\} \qquad (3.10)$$

$$\sigma(nN)_M = (1/4)\{\pi^- p)-\sigma(K^- p)+\sigma(\pi^+ p)-\sigma(K^- n)+\sigma(K^+ p)+\sigma(K^+ N)\} \qquad (3.11)$$

$$\sigma(sN)_B = (1/6)\{\sigma(\Sigma^- p)+\sigma(\Sigma^- n)+\sigma(\Xi^- p)+\sigma(\Xi^- n)-\sigma(pp)-\sigma(pn)\} \quad (3.12)$$

$$\sigma(sN)_m = (1/4)\{\sigma(K^- p)-\sigma(\pi^- p)+\sigma(K^- n)-\sigma(\pi^+ p)+\sigma(K^+ p)+\sigma(K^+ n)\} \quad (3.13)$$

where $\sigma(nN)_B$, $\sigma(nN)_M$, $\sigma(sN)_B$ and $\sigma(sN)_M$ denote the contributions from nonstrange and strange quarks to the isospin averaged baryon-nucleon and meson-nucleon scattering cross sections respectively as calculated from the AQM and the conventional duality assumption of equality of the contributions from strange quarks and antiquarks is used to eliminate antiquark contributions from eq. (3.11) and (3.13),

$$\sigma(sN)_m = \sigma(\bar{s}N)_M \quad (3.14)$$

The AQM predicts the equality of the corresponding quark-nucleon contributions to baryon and meson cross sections. Substituting the relations (3.10)-(3.14) gives two sum rules which can be tested against experimental data:

$$\sigma(nN)_B = \sigma(nN)_M \quad (3.15)$$

$$12.9 \pm 0.01 \text{ mb.} = 11.2 \pm 0.05 \text{ mb.} \quad (3.16)$$

$$\sigma(sN)_B = \sigma(sN)_M \quad (3.17)$$

$$7.7 \pm 0.1 \text{ mb} = 7.75 \pm 0.05 \text{ mb.} \quad (3.18)$$

The experimental data quoted are taken at 100 GeV/c momentum, where there are both new data on hyperon-nucleon cross sections and previous data on the other hadronic cross sections available. The strange sum rule (3.17)-(3.18) is seen to be in excellent agreement with experiment, while there is strong disagreement with the nonstrange sum rule (3.15)-(3.16). The 15% discrepancy is significant and shows that universality holds for the contribution from strange quarks to the hadron-nucleon cross sections, but that the contribution from nonstrange quarks is greater in baryons than in mesons.

The two-component Pomeron model[34] predicts the discrepancy in eq. (3.25)-(3.16) from the quadratic second term on the r.h.s. of eq. (3.8). The availability of the hyperon-nucleon data now allow the two-component Pomeron sum rule (3.7) to be restated and tested at the quark level:

$$\sigma(nN)_B - \sigma(nN)_M = (1/2)\{\sigma(nN)_M - \sigma(sN)_M\} \tag{3.19}$$

$$1.69 \pm 0.05 \text{ mb.} = 1.73 \pm 0.04 \text{ mb.} \tag{3.20}$$

Like eq. (3.7), the sum rule (3.19)-(3.20) relates quantities which vanish respectively in the AQM and in the SU(3) symmetry limit and relates the breaking of additivity between the baryon and meson sectors with the SU(3) breaking in the meson sector.

The indication that strange quark contributions are somehow simpler than nonstrange contributions is a significant and recurrent feature of the data with no explanation from first principles. A pion cloud coupled only to nonstrange quarks is immediately suggested, but has not yet been successfully introduced to give quantitative predictions. This clear signal with no theoretical interpretation offers a challenge to theorists attempting to use QCD for a fundamental explanation of hadron dynamics.

4. THE THEORETICAL BASIS AND PHENOMENOLOGY OF THE OZI RULE

4.1 Introduction - Non-leading Terms, Nonet Symmetry and Exchange Degeneracy

We now return to a more complete discussion of the OZI rule. What exactly is this rule? What does it say? What is its theoretical basis? What are the experimental tests?

A naive formulation of the OZI rule forbids all processes described by disconnected quark-line diagrams. But there is no statement of how much such forbidden processes are suppressed; i.e. whether the suppression factor is 2, 10, 100 or 10^5. There is no prescription for comparison of corresponding allowed for forbidden processes. For example, consider the ratio

$$\frac{\sigma(\pi p \to \phi\pi\pi X)}{\sigma(\pi p \to \phi K\bar{K}X)} = ? \tag{4.1}$$

The $\pi\pi X$ final state is forbidden if X has no strange particles. The $K\bar{K}X$ final state is allowed, but it contains an additional strange quark pair. How does one compare the price paid for producing a pair of strange particles with the price paid for violating OZI? The answer depends on finding the right "non-leading" term responsible for the OZI violating.

The essential theoretical difficulty is the "higher order paradox". A transition forbidden by the OZI rule can proceed as a two-step process in which each of the individual steps is allowed, e.g.

$$\phi \to K^+ + K^- \to \rho + \pi , \tag{4.2}$$

$$f' \to K^+ + K^- \to \pi + \pi , \tag{4.3}$$

$$\pi^- + p \to K^+ + K^- + n \to \phi + n . \tag{4.4}$$

Extensive invetigations of these two-step transitions have shown that the nonet symmetry and exchange degenaracy which appear automatically in the duality description of hadron scattering processes lead to destructive interference and cancellations between contributions from different intermediate states and play a crucial role in suppressing these OZI-violating higher order transitions. This intimate connection between the OZI rule and duality must somehow come out of the QCD description of hadron dynamics, but we have not yet reached the stage of understanding it from first principles.

The role of cancellations, symmetries and symmetry breaking in OZI-forbidden processes is illustrated in the example of the decay of a high-mass charmed D^* meson resonance into a D meson and a kaon pair[40]

$$D^* (c\bar{d}) \to D^+ + K^+ + K^- , \tag{4.5a}$$

$$D^{*+} (c\bar{d}) \to D^+ + K^0 + \bar{K}^0 . \tag{4.5b}$$

The naive OZI rule which forbids all "disconnected quark-line diagrams" forbids the decay mode (4.5a) with a charged kaon pair, but allows the decay mode (4.5b) with a neutral kaon pair,

$$\langle D^+K^+K^-|T|D^{*+}\rangle = \langle(c\bar{d})(u\bar{s}s\bar{u})|T|c\bar{d}\rangle = 0 \;, \tag{4.6a}$$

$$\langle K^0\bar{K}^0D^+|T|D^{*+}\rangle = \langle(c\bar{d}d\bar{s}s\bar{d})|T|c\bar{d}\rangle \neq 0 \;, \tag{4.6b}$$

where T denotes the transition operator for the decay process. The quark lines of the $u\bar{u}$ and $s\bar{s}$ pairs in the charged kaon pair of the decay (4.6a) cannot be connected to the c and $\bar{d}$ quark lines in the D^{*+} and D^+ mesons and the decay is forbidden. But the quark lines of the $d\bar{d}$ pair in the neutral kaon pair of the decay (4.6b) are easily connected with the $\bar{d}$ quark lines in the D^{*+} and D^+ to give an allowed decay.

If the kaon pair is in a state of definite isospin, either 0 or 1, the transition matrix elements (4.6a) and 4.6b) are required to be equal by isospin invariance. Thus, the OZI rule will be violated if the kaon pair is produced via the decay of a resonance with a well-defined isospin, e.g. the f^0 isoscalar or the A_2^0 isovector tensor meson,

$$\langle K^0\bar{K}^0|d|f^0\rangle\langle D^+f^0|p|D^{*+}\rangle = \langle K^+K^-|d|f^0\rangle\langle D^+f^0|p|D^{*+}\rangle, \tag{4.7a}$$

$$\langle K^0\bar{K}^0|d|A_2^0\rangle\langle D^+A_2^0|p|D^{*+}\rangle = -\,\langle K^+K^-|d|A_2^0\rangle\langle D^+A_2^0|p|D^{*+}\rangle, \tag{4.7b}$$

where p and d denote the operators for the production and decay transitions respectively. All these decays are two-step processes, in which each step is clearly allowed by the OZI rule. The contributions from either the f^0 or the A_2^0 intermediate states to the two decays (4.5) are required to be equal by isospin invariance, even though the overall transition (4.6a) involves a disconnected diagram. The OZI violation comes about because the f^0 and A_2 mesons each contain both $u\bar{u}$ and $d\bar{d}$ components and are produced via the $d\bar{d}$ component and decay via the $u\bar{u}$ component.

However, in the nonet symmetry limit, the OZI rule is saved because the f^0 and A_2 contributions to the forbidden reaction (4.5a) exactly cancel. The f^0 and A_2 mesons are exactly degenerate in this limit and their amplitudes satisfy the relations

$$\langle K^0\bar{K}^0|d|f^0\rangle\langle D^+f^0|p|D^{*+}\rangle = \langle K^0\bar{K}^0|d|A_2^0\rangle\langle D^+A_2^0|p|D^{*+}\rangle, \tag{4.8a}$$

$$\langle K^+K^-|d|f^0\rangle\langle D^+f^0|p|D^{*+}\rangle = -\langle K^+K^-|d|A_2^0\rangle\langle D^+A_2^0|p|D^{*+}\rangle. \tag{4.8b}$$

The transition matrix elements for the overall transition are thus given by

$$\langle K^+K^-D^+|T|D^{*+}\rangle = \langle K^+K^-|d|f^0\rangle\langle D^+f^0|p|D^{*+}\rangle + \langle K^+K^-|d|A_2^0\rangle\langle D^+A_2^0|p|D^{*+}\rangle = 0, \tag{4.9a}$$

$$\langle K^0\bar{K}^0D^+|T|D^{*+}\rangle = \langle K^0\bar{K}^0|d|f^0\rangle\langle D^+f^0|p|D^{*+}\rangle + \langle K^0\bar{K}^0|d|A_2^0\rangle\langle D^+A_2^0|p|D^{*+}\rangle$$

$$= 2\langle K^0\bar{K}^0|d|f^0\rangle\langle D^+f^0|p|D^{*+}\rangle. \tag{49.b}$$

The overall transition thus satisfies the naive OZI rule which forbids all disconnected diagrams, even though the forbidden process has contributions from two-step processes in which each step is allowed.

The nonet symmetry (4.8) is essential to provide the cancellations (4.9a) in the forbidden process. The relation between this symmetry limit and the OZI rule is more easily seen in a basis of quarkonium states which each have only a single flavor. Such states can be defined as linear combinations of the f^0 and A_2^0

$$|T_u(u\bar{u})\rangle = \sqrt{\tfrac{1}{2}}\,(|f^0\rangle + |A_2^0\rangle), \tag{4.10a}$$

$$|T_d(d\bar{d})\rangle = \sqrt{\tfrac{1}{2}}\,(|f^0\rangle - |A_2^0\rangle). \tag{4.10b}$$

In this basis the OZI rule gives simple selection rules for the production and decay vertices:

$$\langle K^+K^-|d|T_d\rangle = 0 = \langle D^+T_u|p|D^{*+}\rangle. \tag{4.11a}$$

The selection rule forbidding the overall transition (4.5a) immediately follows:

$$\langle K^+K^-D^+|T|D^{*+}\rangle$$

$$= \langle K^+K^-|d|T_d\rangle\langle D^+T_d|p|D^{*+}\rangle + \langle K^+K^-|d|T_u\rangle\langle D^+T_u|p|D^{*+}\rangle = 0. \qquad (411.b)$$

In this description both terms in eq. (4.11b) vanish individually and no cancellation is required.

This example shows that the OZI rule can hold only in the nonet symmetry limit when the contributing intermediate states can be expressed in a basis of physical quarkonium states with a single flavor. When nonet symmetry is broken, the physical states are always flavor mixtures. The propagators of such intermediate mixed states break the OZI rule. The intermediate state can be produced via one component and decay via another.

However, neither nonet symmetry nor exchange degeneracy are exact in the real world, and in particular there are strong deviations near thresholds. The reactions (4.2)-(4.3) are close enough to the $K\bar{K}$ threshold so that the $K\bar{K}$ channel is open but other channels involving K^* mesons related by exchange degeneracy are closed and cannot provide the necessary cancellations of the amplitudes. These higher order transitions lead to OZI violations which cannot easily be calculated from first principles.

4.2 Phenomenology of Charmonium Decays and Three-meson Photoproduction

A recent analysis[40)] compares and suggests the existence of two degrees of forbiddenness, strongly forbidden and semi-forbidden, corresponding to two distinctly different topologies of disconnected diagrams. We summarize the results of this analysis here.

Table 2 lists some typical examples from charmonium decays and three-meson photoproduction.

Table 2 Experimental values of charmonium decay branching ratios and three-meson photoproduction cross sections

Transitions	OZI Topology	BR(%)	R
$J/\psi \to \omega K\bar{K}$	single-hairpin(S)	0.16±0.10	0.022±0.013
$J/\psi \to \phi\pi^+\pi^-$	single-hairpin(2S)	0.18±0.08	0.024±0.012
$J/\psi \to \phi\pi^+\pi^-$	double-hairpin(S)	0.21±0.09	0.028±0.013
$J/\psi \to \omega\pi\pi$	single-hairpin	0.68±0.19	0.09±0.03
$J/\psi \to \rho\pi$	single-hairpin	1.22±0.12	0.16±0.03
$\psi'(3685) \to K^+K^-\pi^+\pi^-$	single-hairpin(S)	0.16±0.04	0.18±0.05
$\psi'(3685) \to \psi\pi\pi$	ski-track	50±3	56±7
$\psi(3770) \to D\bar{D}$	allowed	100	90000±17000
$\gamma \to \omega\pi^+\pi^-$	allowed		0.93±0.08
$\gamma \to K^{\neq 0}K^{\neq}\pi$	allowed		0.26±0.03
$\gamma \to K^+K^-\rho^0$	allowed		0.13±0.01
$\gamma \to \phi\pi^+\pi^-$	ski-track		0.091±0.017
$\gamma \to \phi\pi^+\pi^-\pi^+\pi^-$	ski-track		0.067±0.013
$\psi \to K^+K^-\omega$	allowed		0.045±0.013
$\psi \to \phi\pi^+\pi^-\pi^0$	single-hairpin(S)		0.020±0.011
$\psi \to K^+K^-V_d$	ski-track(S)		>0.006±0.002
$\psi \to K^+K^-\phi$	allowed(S)		0.005±0.006

The strengths of the various transitions are expressed by a ration R normalized for charmonium decays to the e^+e^- decay of the same initial state and for photoproduction to the cross section for the final state $KK\pi\pi$.

$$R(\psi_i \to X) = \frac{BR(\psi_i \to X)}{BR(\psi_i \to e^+e^-)} , \tag{4.12a}$$

$$R(\gamma \to X) = \frac{\sigma(\gamma \to X)}{\upsilon(\gamma \to KK\pi\pi)} . \tag{4.12b}$$

Allowed decays above the $D\bar{D}$ threshold have values of R of order 10^5. Values of the order 10^{-1}, smaller by a factor of 10^6, are found for all OZI-forbidden charmonium decays except $\psi'(3685)\to\psi\pi\pi$ which is of order 10^2, i.e. three orders of magnitude larger than the other forbidden decays and three orders of magnitude smaller than the allowed decays. The existence of this intermediate-strength decay suggests that some transitions may be only semi-forbidden.

Evidence for such OZI semi-forbiddenness was found in the photoproduction results. The OZI-forbidden $\gamma\to\phi\pi\pi$ was only a factor of three in amplitude smaller than the strongest allowed process measured, $\gamma\to\omega\pi^+\pi^-$ and comparable to the allowed $\gamma\to K^+K^-\rho^0$ and $\gamma\to K^+K^-\omega$, which also involved the creation of a strange quark pair. This first suggested that there was no OZI suppression at all in $\Gamma\to\phi\pi\pi$ and that all the suppression relative to $\gamma\to\omega\pi^+\pi^-$ could be attributed entirely to the creation of a strange quark pair.

A more detailed investigation showed that there was indeed some OZI suppression in addition to the strange quark factor, but that this suppression was much less than in the standard examples of OZI-forbidden processes with light quarks. The data were fit by a phenomenological parameterization which introduced suppression factors for both OZI forbiddenness and strange quark pair creation. A key ingredient in this analysis was the elimination of strange quark suppression factors by analyzing the OZI violation in the two processes $\gamma\to K^+K^-\rho^0$ and $\gamma\to K^+K^-\omega$ which involve hadrons with the same masses and have the same strange quark suppression and kinematic factors. They are predicted to be equal by the OZI rule and are definitely unequal experimentally.

4.3 A Topological Criterion for a Forbiddenness Hierarchy

Ref. [40] defined two different topologies of OZI-forbidden disconnected diagrams:

1. Hairpin diagrams (strongly forbidden), in which the quark lines from a single hadron are disconnected from all other hadrons. Such diagrams remain forbidden under all possible crossing transformations. All the three standard forbidden processes (4.2)-(4.4) and all the strongly forbidden charmonium decays in Table 2 are described by this type of diagram.

2. Ski-track diagrams (semi-forbidden), in which the quark lines from two hadrons are connected together and disconnected from the remaining

part of the diagram, but there is no disconnected single hadron. The term "crossed pomeron" was originally used as an empirical description of the topology of the diagram, but caused confusion because of irrelevant dynamical implications regarding analyticity or crossing symmetry.[41] We avoid this by noting that the two disconnected hadron lines can be ski-tracks, whereas the disconnected line of a hairpin diagram cannot be made by a skier. The reactions

$$\psi'(3685) \to \psi\pi\pi \qquad (4.13a)$$

and

$$\psi + \pi \to \psi'(3685) + \pi \qquad (4.13b)$$

are described by ski-track diagrams. It was conjectured that the suppression factor for such processes might be considerably less than the suppression factor for processes described by hairpin diagrams.

The analysis assumed that hairpin diagrams were forbidden and neglected their contributions, but included the ski-track contributions with a strength proportional to a parameter which was determined by fitting the experimental data. The analysis showed that the contributions of ski-track diagrams was appreciable, and that the suppression factor was much smaller than for contributions from hairpin diagrams.

The transitions in Table 2 are classified by their OZI topology as allowed, single-hairpin, double-hairpin or ski-track processes, with the notation (S) and (2S) added to indicate if one or two strange quark pairs are created in the transition. Processes described by ski-track diagrams are seen to be systematically suppressed by a considerably smaller factor than processes described by hairpin diagrams.

Additional confirmation of semi-forbiddenness is found in the results for $\gamma\to\phi\pi^+\pi^-\pi^0$ and $\gamma\to\phi\pi^+\pi^-\pi^+\pi^-$, although the data in Table 1 are not sufficiently good to draw definite quantitative conclusions. The process $\gamma\to\phi\pi^+\pi^-\pi^0$ must be produced by a hairpin diagram, since the final state has been G-parity and cannot be produced by the odd-G $s\bar{s}$ component of the photon. It seems to be suppressed relative to the states with two of four pions, which can be produced via a ski-track diagram.

The OZI rule remains a puzzle to be unraveled by QCD theory with possible illumination from new experimental data on possible hierarchies of forbiddenness. Investigation of this puzzle may help in our understanding of how hadron structure and dynamics arise from QCD.

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